# Proto-Neutron Star Winds, Magnetar Birth, and Gamma-Ray Bursts

Brian D. Metzger*,†, Todd A. Thompson** and Eliot Quataert*

*Astronomy Department and Theoretical Astrophysics Center, 601 Campbell Hall, Berkeley, CA 94720; bmetzger@astro.berkeley.edu, eliot@astro.berkeley.edu
†Department of Physics, 366 LeConte Hall, University of California, Berkeley, CA 94720
**Department of Astrophysical Sciences, Peyton Hall-Ivy Lane, Princeton University, Princeton, NJ 08544; thomp@astro.princeton.edu

**Abstract.** We begin by reviewing the theory of thermal, neutrino-driven proto-neutron star (PNS) winds. Including the effects of magnetic fields and rotation, we then derive the mass and energy loss from magnetically-driven PNS winds for both relativistic and non-relativistic outflows, including important multi-dimensional considerations. With these simple analytic scalings we argue that proto-magnetars born with $\sim$ millisecond rotation periods produce relativistic winds just a few seconds after core collapse with luminosities, timescales, mass-loading, and internal shock efficiencies favorable for producing long-duration gamma-ray bursts.

**Keywords:** neutron stars, stellar winds, supernovae, gamma ray bursts, magnetic fields
**PACS:** 97.60.Bw, 97.60.Gb; 97.10.Me

## 1. NEUTRINO-DRIVEN PNS WINDS

After a successful core-collapse supernova (SN), a hot proto-neutron star (PNS) cools and deleptonizes, releasing the majority of its gravitational binding energy ($\sim 3\times10^{53}$ ergs) in neutrinos. With initial core temperature $T > 10$ MeV, a PNS is born optically-thick to neutrinos of all flavors because the relevant neutrino-matter cross sections scale as $\sigma_{\nu n} \propto \epsilon_\nu^2 \propto T^2$, where $\epsilon_\nu$ is a typical neutrino energy. Indeed, because neutrinos are trapped, a PNS's neutrino luminosity $L_\nu$ remains substantial and quasi-thermal for a time after bounce $\tau_{\rm KH} \sim 10-100$ s, as roughly verified by the 19 neutrinos detected from SN1987A 20 years ago [1],[2]. Although this Kelvin-Helmholtz (KH) cooling epoch is short compared to the time required for the shock, once successful and moving outward at $\sim 10^4$ km/s, to traverse the progenitor stellar mantle, $\tau_{\rm KH}$ is still significantly longer than the time over which the initial explosion must be successful. While the specific shock launching mechanism is presently unknown, it must occur in a time $t < 1$ s $\ll \tau_{\rm KH}$ after bounce for the PNS to avoid accreting too much matter.

Thus, even after the SN shock has cleared a cavity of relatively low density material around the PNS, $L_\nu$ remains substantial. Detailed PNS cooling calculations [3] show that the electron neutrino(antineutrino) luminosity $L_{\nu_e(\bar{\nu}_e)}$ is $\sim 10^{52}$ erg/s at $t \sim 1$ s and declines as $\propto t^{-1}$ until $t \simeq \tau_{\rm KH}$, after which $L_{\nu_e(\bar{\nu}_e)}$ decreases exponentially as the PNS becomes optically thin. This persistent neutrino flux $F_{\nu_e(\bar{\nu}_e)}$ continues to heat the PNS atmosphere, primarily through electron neutrino(antineutrino)

absorption on nuclei ($\nu_e + n \rightarrow p + e^-$ and $\bar{\nu}_e + p \rightarrow n + e^+$). Because the inverse, pair capture rates dominate the cooling, which declines rapidly with temperature ($\dot{q}^- \propto T^6$) and hence with spherical radius $r$, a region of significant net positive heating ($\dot{q} \equiv \dot{q}^+ - \dot{q}^- > 0$) develops above the neutrinosphere radius $R_\nu$. This heating drives mass-loss from the PNS in the form of a thermally-driven wind [4]. To estimate the dependence of the resultant mass-loss rate ($\dot{M}_{\rm th}$) on the PNS properties explicitly, consider that in steady state the change in gravitational potential required for a unit mass element to escape the PNS ($GM/R_\nu$) must be provided by the total heating it receives accelerating outwards from the PNS surface:

$$\frac{GM}{R_\nu} \approx \int_{R_\nu}^{\infty} \dot{q}\frac{dr}{v_r}, \tag{1}$$

where $M$ is the PNS mass, $v_r$ is the outward wind velocity, and $\dot{q}$ is per unit mass. Because $\dot{q}$ is quickly dominated by heating from neutrino absorption, which scales as $\dot{q}^+ \propto F_{\nu_e}\sigma_{n\nu} \propto L_{\nu_e}\epsilon_{\nu_e}^2/4\pi r^2$, we see that equation (1) implies that

$$\frac{GM}{R_\nu} \propto \frac{L_{\nu_e}\epsilon_{\nu_e}^2}{\dot{M}_{\rm th}}\int_{R_\nu}^{\infty}\rho dr \approx \frac{L_{\nu_e}\epsilon_{\nu_e}^2}{\dot{M}_{\rm th}}\rho_\nu H_\nu, \tag{2}$$

where we have used $\dot{M}_{\rm th} = 4\pi\rho r^2 v_r$ for a spherical wind, $\rho$ is the mass density, $H$ is the PNS's density scale height, $\epsilon_{\nu_e}$ crudely defines a mean electron neutrino or antineutrino energy, and a subscript "$\nu$" denotes evaluation near $R_\nu$. Neglecting rotational support and assuming that the thermal pressure $P$ is dominated by photons and relativistic pairs (which also becomes an excellent approximation as the density plummets abruptly above the PNS surface), we have that $H_\nu \sim P_\nu/\rho_\nu g_\nu \propto T_\nu^4 R_\nu^2/M\rho_\nu$, where $g_\nu \propto M/R_\nu^2$ is the PNS surface gravity and $T_\nu \propto (L_{\nu_e}\epsilon_{\nu_e}^2/R_\nu^2)^{1/6}$ is the PNS surface temperature. $T_\nu$ is set by the balance between heating and cooling at the PNS surface ($T_\nu^6 \propto \dot{q}^- = \dot{q}^+ \propto L_{\nu_e}\epsilon_{\nu_e}^2/R_\nu^2$). Inserting these results into equation (2) and including the correct normalization from the relevant weak cross sections, one finds the expression for $\dot{M}_{\rm th}$ first obtained by ref [4]:

$$\dot{M}_{\rm th} \approx 10^{-4} L_{52}^{5/3}\epsilon_{10}^{10/3} M_{1.4}^{-2} R_{10}^{5/3}\, M_\odot/{\rm s}, \tag{3}$$

where $L_{52} \equiv L_{\nu_e} \times 10^{52}$ erg/s, $\epsilon_{10} \equiv 10\epsilon_{\nu_e}$MeV, $R_\nu \equiv 10 R_{10}$ km, and $M \equiv 1.4 M_{1.4} M_\odot$.

Endowed with an enormous gravitational binding energy and a means, through this neutrino-driven outflow, for communicating a fraction of this energy to the outgoing shock, a newly-born PNS seems capable of affecting the properties of the SN that we observe. However, a purely thermal, neutrino-driven PNS wind is only accelerated to an asymptotic speed of order the surface sound speed: $v_{\rm th}^\infty \sim c_{s,\nu} \approx \sqrt{2kT_\nu/m_p} \approx 0.1 L_{52}^{1/12}\epsilon_{10}^{1/6} R_{10}^{-1/6}$ c. Thus, the efficiency $\eta$ relating wind power $\dot{E}_{\rm th} \approx \dot{M}_{\rm th}(v_{\rm th}^\infty)^2/2$ to total neutrino luminosity ($L_\nu \sim 6L_{\nu_e}$) is quite low:

$$\eta \equiv \frac{\dot{E}_{\rm th}}{L_\nu} \sim 10^{-5} L_{52}^{5/6}\epsilon_{10}^{11/3} R_{10}^{4/3} M_{1.4}^{-2}. \tag{4}$$

In particular, although neutrino energy deposited in a similar manner may be responsible for initiating the SN explosion itself at early times (i.e., the neutrino SN mechanism [5]), $\eta$ drops rapidly as the PNS cools. Quasi-spherical winds of this type are therefore not expected to affect the SN's nucleosynthesis or morphology (although the wind itself is considered a promising $r$-process source [4]).

## 2. MAGNETICALLY-DRIVEN PNS WINDS

Some PNSs may possess a more readily extractable form of energy in rotation. A PNS born with a period $P = P_{\rm ms}$ ms is endowed with a rotational energy $E_{\rm rot} \simeq 2\times 10^{52} P_{\rm ms}^{-2} R_{10}^2 M_{1.4}$ ergs, which, for $P < 4$ ms, exceeds the energy of a typical SN shock ($\sim 10^{51}$ ergs). Given a mass loss rate $\dot{M}$ and torquing lever arm $\omega_\tau$, a wind extracts angular momentum $J$ from the PNS at a rate $\dot{J} \simeq \Omega\omega_\tau^2\dot{M}$, where $\Omega = 2\pi/P$ is the PNS rotation rate. With the PNS's radius $R_\nu$ as a lever arm and the modest thermally-driven mass-loss rate given by equation (3), the timescale for removal of the PNS's rotational energy, $\tau_J \equiv J/\dot{J} \sim MR_\nu^2/\dot{M}\omega_\tau^2 \sim M/\dot{M}_{\rm th}$, is much longer than $\tau_{\rm KH}$. However, if the PNS is rapidly rotating and possesses a dynamically-important poloidal magnetic field $B_p$ (through either flux-freezing or generated via dynamo action [6]), then both $\dot{M}$ and $\omega_\tau$ can be substantially increased; this reduces $\tau_J$, allowing efficient extraction of $E_{\rm rot}$.

For magnetized winds $\omega_\tau$ is the Alfvén radius $\omega_A$, defined as the cylindrical radius where $\rho v_r^2/2$ first exceeds $B_p^2/8\pi$ [7]. The magnetosphere of a PNS is most likely dominated by its dipole component, with a total (positive-definite) surface magnetic flux given by $\Phi_{\rm B} = 2\pi B_\nu R_\nu^2$, where $B_\nu$ is the polar surface field. To estimate $\omega_A$ for magnetized PNS outflows recognize that mass and angular momentum are primarily extracted from a PNS along open magnetic flux. For an axisymmetric dipole rotator this represents only a fraction $\approx 2(\pi\theta_{\rm LCFL}^2)/4\pi \simeq R_\nu/2\omega_{\rm Y}$ of $\Phi_{\rm B}$, where $\theta_{\rm LCFL} \approx \sqrt{R_\nu/\omega_{\rm Y}}$ is the latitude (measured from the pole) at the PNS surface of the last closed field line (LCFL), $\omega_{\rm Y}$ is the radius where the LCFL intersects the equator (the "Y point"), and we have assumed that $\omega_{\rm Y} \gg R_\nu$ ($\theta_{\rm LCFL} \ll 1$). Plasma necessarily threads a PNS's closed magnetosphere and cannot be forced to corotate superluminally; thus $\omega_{\rm Y}$ cannot exceed the light cylinder radius $\omega_{\rm L} \equiv c/\Omega = 48P_{\rm ms}$ km, making it useful to write the PNS magnetosphere's total open magnetic flux as $\Phi_{\rm B,open} \approx \pi B_\nu R_\nu^2 (R_\nu/\omega_{\rm L})(\omega_{\rm Y}/\omega_{\rm L})^{-1}$. Now, the overall *latitudinal structure* of a PNS magnetosphere (i.e., the allocation of open and closed magnetic flux, and the value of $\omega_{\rm Y}/\omega_{\rm L}$) is primarily dominated by the dipolar closed zone. However, recent numerical simulations [8] show that where the field is open it behaves as a "split monopole". In this case the poloidal field scales as $B_p \sim \Phi_{\rm B,open}/r^2 \approx 0.2 B_\nu P_{\rm ms}^{-1} R_{10} (\omega_{\rm Y}/\omega_{\rm L})^{-1}(R_\nu/r)^2$, rather than the dipole scaling $\propto (R_\nu/r)^3$. The constant of proportionality is chosen to assure that $B_p(R_\nu) \to B_\nu$ in the limit of vanishing closed zone ($\omega_{\rm L}, \omega_{\rm Y} \to R_\nu$) and is in agreement with numerical results (see eq. [28] of ref [8]).

## 2.1. Non-Relativistic Winds and Asymmetric Supernovae

Non-relativistic (NR) magnetically-driven winds reach an equipartition between kinetic and magnetic energy outside $\omega_A$ such that the kinetic energy flux at $\omega_A$ $(\dot{M}v_r(\omega_A)^2/2)$ carries a sizeable fraction of the rotational energy loss extracted by the wind's surface torque $\dot{E}_{\rm rot} = \dot{J}\Omega = \dot{M}\Omega^2\omega_A^2$; thus, we have that $v_r(\omega_A) \sim \Omega\omega_A$. Combining this with the modified monopole scaling for $B_p$ motivated above and mass conservation $\dot{M}_\Omega \equiv \rho r^2 v_r$ ($\dot{M}_\Omega$ is the mass flux *per solid angle*) we find that:

$$\omega_{\rm A}/R_\nu \simeq B_{15}^{2/3}P_{\rm ms}^{-2/3}\dot{M}_{\Omega,-4}^{-1/3}R_{10}^{4/3}(\omega_{\rm Y}/\omega_{\rm L})^{-1}, \tag{5}$$

where $\dot{M}_\Omega \equiv \dot{M}_{\Omega,-4}\times 10^{-4}M_\odot{\rm s}^{-1}{\rm sr}^{-1}$, $B_\nu \equiv B_{15}\times 10^{15}$ G, and we have concentrated on the open magnetic flux that emerges nearest the closed zone (polar latitude $\approx \theta_{\rm LCFL}$) and which thereby dominates the spin-down torque.

From equation (5) we see that winds from rapidly rotating PNSs with surface magnetic fields typical of Galactic "magnetars" ($B_\nu \sim 10^{14} - 10^{15}$ G) possess enhanced lever arms for extracting rotational energy [9]. Furthermore, their total outflow power $\dot{E}_{\rm mag}^{\rm NR} \approx \dot{E}_{\rm rot} \approx 2\pi\theta_{\rm LCFL}^2\dot{M}_\Omega\Omega^2\omega_A^2 \approx 10^{49}B_{15}^{4/3}P_{\rm ms}^{-13/3}\dot{M}_{\Omega,-4}^{1/3}R_{10}^{17/3}(\omega_{\rm Y}/\omega_{\rm L})^{-3}$ ergs/s dominates thermal acceleration ($\dot{E}_{\rm mag}^{\rm NR} > \dot{E}_{\rm th}$) for $B_{15} > 0.4P_{\rm ms}^{13/4}L_{52}^{23/24}\epsilon_{10}^{23/12}R_{10}^{-11/3}M_{1.4}^{-1}(\omega_{\rm Y}/\omega_{\rm L})^{9/4}$. This condition becomes easier to satisfy as the PNS cools, allowing magnetized winds to dominate later stages of the KH epoch for PNSs with even relatively modest $B_\nu$ and $\Omega$. NR magnetically-driven winds, in addition to being more powerful than spherical, thermally-driven outflows, are efficiently hoop-stress collimated along the PNS rotation axis [8]. The power they deposit along the poles may produce asymmetry in SN ejecta distinct from the shock-launching process itself.

Strong magnetic fields and rapid rotation can also increase the outflow's power through enhanced mass-loss because $\dot{E}_{\rm mag}^{\rm NR} \propto \dot{M}_\Omega^{1/3}$. When the PNS's hydrostatic atmosphere is forced to co-rotate to the outflow's sonic radius $\omega_{\rm s} = (GM\sin[\theta_{\rm LCFL}]/\Omega^2)^{1/3}$ then $\dot{M}_\Omega$ is enhanced by a factor $\phi_{\rm cf} \sim \exp[(v_{\phi,\nu}/c_{s,\nu})^2]$ over $\dot{M}_{\rm th}/4\pi$ due to centrifugal ("cf") slinging [9], where $v_{\phi,\nu} \approx R_\nu\Omega\sin[\theta_{\rm LCFL}] \approx R_\nu\Omega\sqrt{R_\nu/\omega_{\rm Y}}$ is the PNS rotation speed at the base of the open flux. Using our estimate for $c_{s,\nu}$ from § 1, we see that enhanced mass loss becomes important for $P_{\rm ms} < P_{\rm cf,ms} \equiv L_{52}^{-1/18}\epsilon_{10}^{-1/9}R_{10}^{10/9}(\omega_{\rm Y}/\omega_{\rm L})^{-1/3}$ (i.e., only for PNSs with considerable rotational energy $E_{\rm rot} > 10^{52}$ ergs). Fully enhanced mass loss ($\dot{M}_\Omega = \dot{M}_{\rm th}\phi_{\rm cf}/4\pi$) requires $\omega_A > \omega_{\rm s}$, which in turn requires that $B_{15} > B_{\rm cf,15} \equiv P_{\rm ms}^{7/4}R_{10}^{-13/4}\dot{M}_{\Omega,-4}^{1/2}M_{1.4}^{1/2}(\omega_{\rm Y}/\omega_{\rm L})^{5/4} \simeq 0.3P_{\rm ms}^{7/4}L_{52}^{5/6}\epsilon_{10}^{5/3}M_{1.4}^{-1/2}R_{10}^{-29/12}\exp[0.5(P/P_{\rm cf})^{-3}](\omega_{\rm Y}/\omega_{\rm L})^{5/4}$, where we have taken $\dot{M}_{\rm th}$ from § 1. For cases with $B_\nu < B_{\rm cf}$ but $P < P_{\rm cf}$, $\dot{M}_\Omega$ lies somewhere between $\dot{M}_{\rm th}/4\pi$ and $\phi_{\rm cf}\dot{M}_{\rm th}/4\pi$ (see [10] for numerical results). Millisecond proto-magnetars generally attain $\phi_{\rm cf}$, except perhaps at early times when the PNS is quite hot.

### 2.2. Relativistic Winds and Gamma-Ray Bursts

As the PNS cools, eventually $\omega_A \rightarrow \omega_{\rm L}$ and the PNS outflow becomes relativistic (REL). This transition occurs after $\tau_{\rm KH}$ for most PNSs (they become pulsars), but rapidly rotating proto-magnetar winds become relativistic during the KH epoch itself. Similar to normal pulsars, PNSs of this type lose energy at the force-free, "vacuum dipole" rate: $\dot{E}_{\rm mag}^{\rm REL} \approx 6\times 10^{49} B_{15}^2 P_{\rm ms}^{-4} R_{10}^6 (\omega_{\rm Y}/\omega_{\rm L})^{-2}$ ergs/s (again modulo corrections for excess open magnetic flux $\dot{E}_{\rm mag}^{\rm REL} \propto \Phi_{\rm B,open}^2 \propto (\omega_{\rm Y}/\omega_{\rm L})^{-2}$ [8]), which gives a familiar spin-down timescale $\tau_{\rm J} = E_{\rm rot}/\dot{E}_{\rm mag}^{\rm REL} \approx 300 B_{15}^{-2} P_{\rm ms}^2 R_{10}^{-4} M_{1.4} (\omega_{\rm Y}/\omega_{\rm L})^2$ s. On the other hand, the mass loading on a PNS's open magnetic flux is set by neutrino heating, a process totally different from the way that matter is extracted from a normal pulsar's surface. In fact, a proto-magnetar outflow's energy-to-mass ratio $\sigma$ is given by

$$\sigma \approx \frac{\dot{E}_{\rm mag}^{\rm REL}}{2\pi\theta_{\rm LCFL}^2 \dot{M}_\Omega c^2} \approx 3 B_{15}^2 P_{\rm ms}^{-3} L_{52}^{-5/3} \epsilon_{10}^{-10/3} R_{10}^{10/3} M_{1.4}^2 \exp\left[-\left(\frac{P}{P_{\rm cf}}\right)^{-3}\right]\left(\frac{\omega_{\rm Y}}{\omega_{\rm L}}\right)^{-1} \quad (6)$$

From equation (6) we see that because a PNS's mass-loss rate drops so precipitously as it cools, $\sigma \propto L_{\nu_e}^{-5/3} \epsilon_{\nu_e}^{-10/3}$ rises rapidly with time, easily reaching $\sim 10-1000$ during the KH epoch for typical magnetar parameters [9],[10]. Detailed evolution calculations indicate that $E_{\rm rot}$ is extracted roughly uniformly in $\log(\sigma)$ [10].

To conclude with a concrete example, consider a proto-magnetar with $B_\nu = 10^{16}$ G and $P_{\rm ms} = 3$ at $t = 10$ seconds after core collapse. From the cooling calculations of ref [3] we have $L_{52}(10\ {\rm s}) \approx 0.1$ and $\epsilon_{10}(10\ {\rm s}) \approx 1$ (see Figs. [14] and [18]) and so, under the conservative estimate that $\omega_{\rm Y} = \omega_{\rm L}$, equation (6) gives $\sigma \approx 500$. Because $\sigma$ represents the potential Lorentz factor of the outflow (assuming efficient conversion of magnetic to kinetic energy), we observe that millisecond proto-magnetar birth provides the right mass-loading to explain gamma-ray bursts (GRBs). Further, the power at $t = 10$ s is still $\dot{E}_{\rm mag}^{\rm REL} \approx 10^{50}$ erg/s with a spin-down time $\tau_{\rm J} \approx 30$ s, both reasonable values to explain typical luminosities and durations, respectively, of long-duration GRBs. Lastly, because $\sigma$ rises so rapidly with time as the PNS cools, in the context of GRB internal shock models a cooling proto-magnetar outflow's kinetic-to-$\gamma$-ray efficiency can be quite high; our calculations indicate that values of $10-50\%$ are plausible. We conclude that magnetar birth accompanied by rapid rotation (but requiring less angular momentum than collapsar models) represents a viable long-duration GRB central engine.

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